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Stable Superstring Relics and Ultrahigh Energy Cosmic Rays

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Abstract

One of the most intriguing experimental results of recent years is the observation of Ultrahigh Energy Cosmic Rays (UHECRs) above the GZK cutoff. Plausible candidates for the UHECR primaries are the decay products of a meta–stable matter state with mass of order $O(10^{12-15}$ GeV), which simultaneously is a good cold dark matter candidate. We study possible meta–stable matter states that arise from Wilson line breaking of GUT symmetries in semi-realistic heterotic string models. In the models that we study the exotic matter states can be classified according to patterns of $SO(10)$ symmetry breaking. We show that cryptons, which are states that carry fractional electric charge $\pm 1/2$, and are confined by a hidden gauge group cannot produce viable dark matter. This is due to the fact that, in addition to the lightest neutral bound state, cryptons give rise to meta–stable charged bound states. However, these states may still account for the UHECR events. We argue that the uniton, which is an exotic Standard Model quark but carries "fractional" $U(1)_{Z'}$ charge, as well as the singleton, which is a Standard Model singlet with "fractional" $U(1)_{Z'}$ charge, do provide viable dark matter candidates and can at the same time explain the observed UHECR events.

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1 Introduction

One of the interesting experimental observations of recent years is the detection of Ultrahigh Energy Cosmic Rays [1], whose observed energy exceed the Greisen–Zatsepin–Kuzmin (GZK) cutoff [2]. There are apparently no astrophysical sources in the local neighbourhood that can account for the events. The shower profile of the highest energy events is consistent with identification of the primary particle as a hadron but not as a photon or a neutrino. The ultrahigh energy events observed in the air shower arrays have muonic composition indicative of hadrons. The problem, however, is that the propagation of hadrons over astrophysical distances is affected by the existence of the cosmic background radiation, resulting in the GZK cutoff on the maximum energy of cosmic ray nucleons $E_{\rm GZK} \leq 10^{20}$ eV. Similarly, photons of such high energies have a mean free path of less than 10Mpc due to scattering from the cosmic background radiation and radio photons. Thus, unless the primary is a neutrino, the sources must be nearby. On the other hand, the primary cannot be a neutrino because the neutrino interacts very weakly in the atmosphere. A neutrino primary would imply that the depths of first scattering would be uniformly distributed in column density, which is contrary to the observations.

The difficulty in finding conventional explanations for UHECR opens the door for innovative approaches. One of the most elegant possibilities is that the UHECR originate from the decay of long–lived super–heavy relics, with mass of the order of 10^{12-15} GeV [3]. In this case the primaries for the observed UHECR would originate from decays in our galactic halo, and the GZK bound would not apply. Furthermore, the profile of the primary UHECR indicates that the heavy particle should decay into electrically charged or strongly interacting particles.

From the particle physics perspective two questions are of interest. The first is the stabilization mechanism which produces a super–heavy state with a lifetime of the order of the universe, while still allowing it to decay and account for the observed UHECR events. The second is how the required mass scale can arise naturally. In a field theoretic model addressing both of these questions amounts to fixing the required parameters by hand. It is therefore of further interest to study the plausible states that may arise from string theory, and whether such states can provide candidates for the UHECR events.

One particular string theory candidate that has been proposed previously is the 'crypton' [4, 5, 6], in the context of the flipped $SU(5)$ free fermionic string model [7]. The 'crypton' is a state that carries fractional electric charge $\pm 1/2$ and transforms under a non–Abelian hidden gauge group, which in the case of the flipped $SU(5)$ "revamped"string model is $SU(4)$. The fractionally charged states are confined and produce integrally charged hadrons of the hidden sector gauge group. The lightest hidden hadron is expected to be neutral with the heavier modes split due to their electromagnetic interactions. A priori, therefore, the 'crypton' is an appealing CDM candidate, which is meta–stable because of the fractional electric charge of the con-

stituent ‘quarks’. This implies that the decay of the exotic hadrons can be generated only by highly suppressed non–renormalizable operators. Effectively, therefore, the events that generate the UHECR are produced by annihilation of the ‘cryptons’ in the confining hidden hadrons. Moreover, the mass scale of the hidden hadrons is fixed by the hidden sector gauge dynamics. Therefore, in the same way that the colour $SU(3)_C$ hadronic dynamics are fixed by the boundary conditions at the Planck scale and the $SU(3)_C$ matter content, the hidden hadron dynamics are set by the same initial conditions and by the hidden sector matter content.

However, we argue here that the ‘crypton’ cannot in fact provide an appealing candidate for CDM. The reason is again its fractional electric charge. In addition to producing an electrically neutral hadron, which is expected to be the lightest hidden hadron, integrally charged hadrons are produced as well. The same reasoning that explains the meta–stability of the lightest neutral hadron also implies that, generically, the hidden, electrically charged, hadrons are semi–stable as well. Namely, the constituent cryptons carry fractional electric charge $\pm 1/2$. Suppose we have an integrally charged tetron which is composed of three $+1/2$, and one $-1/2$, constituents. In order to convert the charged tetron to a neutral one, a $+1/2$ constituent has to convert into a $-1/2$. However, since the cryptons are singlets of $SU(2)_L$ this transition can only proceed via heavy GUT, or string, modes which carry electric charge $+1$. Therefore, the transition from the charged hadrons to the lightest neutral hadron can only proceed by operators which are suppressed by the GUT, or string, unification scales. Effectively, similar to the decay of the lightest neutral hadron, this transition can be generated only by non–renormalizable operators. Although one cannot rule out the possibility that in specific models the decay of the neutral hadrons will arise from operators that are suppressed relative to those that induce the charged hidden hadron decays, generically, we expect both lifetimes to be of the same order of magnitude. Therefore, in addition to the lightest neutral hadron that could account for the UHECR, the ‘crypton’ also gives rise to long–lived charged hadrons, whose number densities are severely constrained. We also study the effect of intermediate scale cryptons on the renormalization group running of the Standard Model gauge couplings and show that coupling unification necessitates the existence of additional colour and electroweak states.

These arguments therefore prompt us to examine whether realistic string models may still produce well motivated candidates to account for the UHECR. We argue that the answer is affirmative. We study the different types of exotic states that arise in semi–realistic string models and their viability as candidates for producing the UHECR. We first discuss the classification of the various states in the string models and their properties. There are several distinct categories of string states that may produce semi–stable matter and we elaborate on the various cases. One general distinction in the string models is between matter states that arise from the untwisted and twisted orbifold sectors, which do not break the GUT symmetry, versus those which arise from the Wilsonian sectors, and which do break the GUT symmetry. The

former sector gives rise to Standard Model states which preserve the GUT structure, whereas states from the Wilsonian sector do not fit into multiplets of the original GUT symmetry and are meta–stable due to discrete symmetries. Another category of states are hidden sector glueballs that were suggested as candidates for Self–Interacting Dark Matter (SIDM) [8]. The viable states, that may give rise to the UHECR events, should possess two important properties. First, they should transform under a non–Abelian hidden sector gauge group that confines at $O(10^{11-13})$ GeV. Second, we argue that the most appealing candidates should be Standard Model singlets. Such states can arise from the "Wilsonian" sector and carry a non–$SO(10)$ charge under the $U(1)_{Z'}$ which is embedded in $SO(10)$, or they can arise from hidden sector glueballs.

2 Quasi–stable matter in realistic string models

In this section we discuss the different classes of exotic states in the string models that may produce UHECR candidates. For concreteness we study these questions in the context of the realistic free fermionic heterotic string models. A general remark on the different string theory constructions is that the heterotic string allows for the embedding of the Standard Model states in $SO(10)$ multiplets, whereas type I constructions only permit constructions that have the generic structure which is a product of $U(n)$ groups. The heterotic string is therefore the only perturbative string theory which is compatible with the orthodox unification scenario, which is highly motivated by the Standard Model multiplet structure and the MSSM gauge coupling unification [9].

Heterotic string models generically give rise to exotic matter states which arise because of the breaking of the non–Abelian unifying gauge symmetry, G, by Wilson–lines [10, 11, 12]. The breaking of the gauge symmetries by Wilson lines results in massless states that do not fit into multiplets of the original unbroken gauge symmetry. This is an important property as it may result in conserved quantum numbers that will indicate the stability of these so-called "Wilsonian" states. The simplest example of this phenomenon is the existence of states with fractional electric charge in the massless spectrum of superstring models [11, 4, 13]. Such states are stable due to electric charge conservation. As there exist strong constraints on their masses and abundance, states with fractional electric charge must be diluted away or be extremely massive. The same "Wilson line" breaking mechanism, which produces matter with fractional electric charge, is also responsible for the existence of states which carry the "standard" charges under the Standard Model gauge group but which carry fractional charges under a different subgroup of the unifying gauge group. For example, if the group G is $SO(10)$ then the "Wilsonian" states may carry non–standard charges under the $U(1)_{Z'}$ symmetry, which is embedded in $SO(10)$ and is orthogonal to $U(1)_Y$. Such states can therefore be long–lived if the $U(1)_{Z'}$ gauge

symmetry remains unbroken down to low energies, or if some residual local discrete symmetry is left unbroken after $U(1)_{Z'}$ symmetry breaking. What is noted is that the heterotic–string construction itself, generically embodies the mechanism that results in semi–stable heavy matter, which in turn may produce viable candidates for the UHECR events. The existence of heavy stable "Wilsonian" matter can therefore be argued to be a "smoking gun" of heterotic string unification.

The realistic models in the free fermionic formulation are generated by a basis of boundary condition vectors for all world–sheet fermions [7, 14, 15, 16, 17, 18, 19]. The basis is constructed in two stages. The first stage consists of the NAHE set, $\{\mathbf{1}, S, b_1, b_2, b_3\}$ [20]. At the level of the NAHE set the gauge group is $SO(10) \times SO(6)^3 \times E_8$, with 48 generations and $N = 1$ supersymmetry. The NAHE set correspond to $Z_2 \times Z_2$ orbifold compactification with non–trivial background fields [21]. The Neveu–Schwarz sector corresponds to the untwisted sector, and the sectors b_1, b_2 and b_3 to the three twisted sectors of the $Z_2 \times Z_2$ orbifold model. In addition to the gravity and gauge multiplets, the Neveu–Schwarz sector produces six multiplets in the 10 representation of $SO(10)$, and several $SO(10)$ singlets transforming under the flavour $SO(6)^3$ symmetries. The sectors b_1, b_2 and b_3 produce the Standard Model matter fields that are embedded in the spinorial 16 representations of $SO(10)$. At the level of the NAHE set all the states in the free fermionic models fall into representations of $SO(10)$, or are $SO(10)$ singlets. Furthermore, at this stage the hidden E_8 is unbroken, hidden matter does not arise, and the models do not provide any candidates for the UHECR events.

The second stage of the basis construction consists of adding to the NAHE set three basis vectors, which correspond to "Wilson lines" in the orbifold formulation. The additional vectors reduce the number of generations to three, one from each sector b_1, b_2 and b_3, and break the gauge symmetries of the NAHE set. The $SO(10)$ symmetry is broken to one of its subgroups $SU(5) \times U(1)$ [7], $SO(6) \times SO(4)$ [15], $SU(3) \times U(1)_{B-L} \times SU(2)_L \times SU(2)_R$ [19], or $SU(3) \times SU(2) \times U(1)_{B-L} \times U(1)_{T_{3_R}}$ [17]. At the same time the hidden E_8 symmetry is broken to one of its subgroups.

The choice of the subgroup of $SO(10)$ that is left unbroken at the string scale determines what kind of exotic matter states can appear in a given string model, as well as affecting the final subgroup of E_8 which remains unbroken by the choice of boundary condition basis vector and GSO phases. Since the superstring derived standard–like models contain the symmetry breaking sectors that arise also in the other models, their massless spectra admits also the exotic representations that can appear in these models. We therefore focus our discussion on the superstring standard–like models.

In the superstring standard–like models, the observable gauge group after application of the generalized GSO projections is $SU(3)_C \times U(1)_C \times SU(2)_L \times U(1)_L \times U(1)^3 \times U(1)^n$. The electromagnetic charge is given by

$$U(1)_{\mathrm{e.m.}} = T_{3_L} + U(1)_Y, \tag{1}$$

where T_{3_L} is the diagonal generator of $SU(2)_L$, and $U(1)_Y$ is the weak hypercharge.

The weak hypercharge is given by*

$$U(1)_Y = \frac{1}{3}U(1)_C + \frac{1}{2}U(1)_L \qquad (2)$$

and the orthogonal combination is given by

$$U(1)_{Z'} = U(1)_C - U(1)_L. \qquad (3)$$

The massless spectrum of the standard–like models contains three chiral generations, each consisting of a 16 of $SO(10)$, decomposed under the final $SO(10)$ subgroup as

$$\begin{aligned}
e_L^c &\equiv [(1, \ \frac{3}{2}); (1, \ 1)]_{(\ 1\ ,\ 1/2\ ,\ 1)} \ ; \ u_L^c \equiv [(\bar{3}, -\frac{1}{2}); (1, -1)]_{(-2/3,1/2,-2/3)}; && (4)\\
d_L^c &\equiv [(\bar{3}, -\frac{1}{2}); (1, \ 1)]_{(1/3,-3/2,1/3)} \ ; \ Q \equiv [(3, \frac{1}{2}); (2, \ 0)]_{(1/6,1/2,(2/3,-1/3))}; && (5)\\
N_L^c &\equiv [(1, \ \frac{3}{2}); (1, -1)]_{(\ 0\ ,\ 5/2\ ,\ 0)} \ ; \ L \equiv [(1, -\frac{3}{2}); (2, \ 0)]_{(-1/2,-3/2,(0,1))}, && (6)
\end{aligned}$$

where we have used the notation

$$[(SU(3)_C \times U(1)_C); (SU(2)_L \times U(1)_L)]_{(Q_Y, Q_{Z'}, Q_{\rm e.m.})}, \qquad (7)$$

and have written the electric charge of the two components for the doublets.

The matter states from the NS sector and the sectors b_1, b_2 and b_3 transform only under the observable gauge group. In the realistic free fermionic models, there is typically one additional sector that produces matter states transforming only under the observable gauge group [17]. These states complete the representations that we identify with possible representations of the Standard Model. In addition to the Standard Model states, semi–realistic superstring models may contain additional multiplets, in the 16 and $\overline{16}$ representation of $SO(10)$, in the vectorial 10 representation of $SO(10)$, or the 27 and $\overline{27}$ of E_6. Such states can pair up to form super-massive states. They can mix with, and decay into, the Standard Model representation unless some additional symmetry, which forbids their decay, is imposed. For example, in the flipped $SU(5)$ superstring models [7], two of the additional vectors which extend the NAHE set produce an additional 16 and $\overline{16}$ representation of $SO(10)$. These states are used in the flipped $SU(5)$ model to break the $SU(5) \times U(1)$ symmetry to $SU(3) \times SU(2) \times U(1)$.

In addition to the states mentioned above transforming solely under the observable gauge group, the sectors $b_j + 2\gamma$ produce matter states that fall into the 16 representation of the hidden $SO(16)$ gauge group decomposed under the final hidden gauge group. The states from the sectors $b_j + 2\gamma$ are $SO(10)$ singlets, but are

* Note that we could have instead defined the weak hypercharge to be $U(1)_Y = \frac{1}{3}U(1)_C - \frac{1}{2}U(1)_L$. This amounts to the same redefinition of fields between the straight and flipped $SU(5)$. In this paper we will use the definition in Eq. 2.

charged under the flavour $U(1)$ symmetries. All the states above fit into standard representations of the grand unified group which may be, for example, $SO(10)$ or E_6, or are singlets of these groups. They carry the standard charges under the Standard Model gauge group or of its GUT extensions.

The superstring models contain additional states that cannot fit into multiplets of the original $SO(10)$ unifying gauge group. They result from the breaking of the $SO(10)$ gauge group at the string level via the boundary condition assignment. The exotic states in the realistic free fermionic models appear in vector–like representations and can acquire a large mass. Next we enumerate the exotic states that appear in free fermionic models. The states are classified according to the unbroken $SO(10)$ subgroup in each sector [12].

From the $\underline{SO(6)\times SO(4)}$ type sectors we obtain the following exotic states.

- Colour triplets : $\quad [(3,\frac{1}{2});(1,0)]_{(1/6,1/2,1/6)} \quad ; \quad [(\bar{3},-\frac{1}{2});(1,0)]_{(-1/6,-1/2,-1/6)}$
- Electroweak doublets : $\quad [(1,0);(2,0)]_{(0,0,\pm1/2)}$
- Fractionally charged $SU(3)_C\times SU(2)_L$ singlets :

$$[(1,0);(1,\pm1)]_{(\pm1/2,\mp1/2,\pm1/2)} \quad ; \quad [(1,\pm3/2);(1,0)]_{(\pm1/2,\pm1/2,\pm1/2)} \tag{8}$$

The colour triplets bind with light quarks to form mesons and baryons with fractional electric charges $\pm1/2$ and $\pm3/2$. The $\underline{SO(6)\times SO(4)}$ type states can appear in the Pati–Salam type models [15] or in Standard–like models [17].

From sectors which break the $SO(10)$ symmetry into $\underline{SU(5)\times U(1)}$ we obtain exotic states with fractional electric charge $\pm1/2$

- Fractionally charged $SU(3)_C\times SU(2)_L$ singlets :

$$[(1,\pm3/4);(1,\pm1/2)]_{(\pm1/2,\pm1/4,\pm1/2)} \tag{9}$$

In general the fractionally charged states may transform under a non–Abelian hidden gauge group in which case the fractionally charged states may be confined. For example, in the "revamped" flipped $SU(5)$ model [7] the states with fractional charge $\pm1/2$ transform as 4 and $\bar{4}$ of the hidden $SU(4)$ gauge group. The states with the charges in eq. (9) are called the "cryptons" and may form good dark matter candidates [4] if the lightest confined state is electrically neutral. In the "revamped" flipped $SU(5)$ model it has been argued that the lightest state is the "tetron", which contains four fundamental constituents. In other models, states with the charges of eq. (9) may be singlets of all the non–Abelian group factors.

Finally in the superstring derived standard–like models we may obtain exotic states from sectors which are combinations of the $\underline{SO(6)\times SO(4)}$ breaking vectors and $\underline{SU(5)\times U(1)}$ breaking vectors. These states therefore arise only in the $\underline{SU(3)\times SU(2)\times U(1)^2}$ type models. These states then carry the standard charges

under the Standard Model gauge group but carry fractional charges under the $U(1)_{Z'}$ gauge group. The following exotic states are obtained:

- colour triplets : $$[(3,\frac{1}{4});(1,\frac{1}{2})]_{(-1/3,-1/4,-1/3)} \quad ; \quad [(\bar{3},-\frac{1}{4});(1,\frac{1}{2})]_{(1/3,1/4,1/3)} \quad (10)$$

Due to its potential role in string gauge coupling unification [24], this state is referred to as “the uniton” [12]. The uniton forms bound states with the lightest up and down quarks and gives rise to ultra–heavy mesons. In ref. [12] it has been shown that the lightest meson can be the electrically neutral state.

- electroweak doublets : $[(1,\pm\frac{3}{4});(2,\pm\frac{1}{2})]_{(\pm 1/2,\pm 1/4,(1,0);(0,-1))}$

Unlike the previous electroweak doublets, these electroweak doublets carry the regular charges under the standard model gauge group but carry “fractional” charge under the $U(1)_{Z'}$ symmetry. Finally, in the superstring derived standard–like models we also obtain states which are Standard Model singlets but carry “fractional” charges under the $U(1)_{Z'}$ symmetry.

- Standard model singlets with “fractional” $U(1)_{Z'}$ charge :

$$[(1,\pm\frac{3}{4});(1,\mp\frac{1}{2})]_{(0,\pm 5/4,0)} \tag{11}$$

These states may transform under a non–Abelian hidden gauge group or may be singlets of all the non–Abelian group factors. This type of Standard Model singlet appears in all the known free fermionic standard–like models. We refer to this state as the “singleton”.

There are several important issues to examine with regard to the exotic states. Since some of these states carry fractional charges, it is desirable to make them sufficiently heavy or sufficiently rare. A priori, in a generic string model, it is not at all guaranteed that the states with fractional electric charge can be decoupled or confined [22]. Therefore, their presence imposes an highly non–trivial constraint on potentially viable string vacua. In the NAHE–based free fermionic models, all the exotic matter states appear in vector–like representations. They can therefore obtain mass terms from renormalizable or higher order terms in the superpotential. We must then study the renormalizable and nonrenormalizable superpotential in the specific models. The cubic level and higher order terms in the superpotential are extracted by applying the rules of ref. [23]. The problem of fractionally charged states was investigated in ref. [13, 18] for the model of ref. [14]. By examining the fractionally charged states and the trilinear superpotential, it was shown that all the fractionally charged states receive a Planck scale mass by giving a VEV to a set of $SO(10)$ singlets in the massless string spectrum. Therefore, all the fractionally charged states can decouple from the remaining light spectrum. The second issue that must be examined

with regard to the exotic "Wilsonian" matter is the interactions with the Standard Model states. The fractional charges of the exotic states under the unbroken $U(1)$ generators of the $SO(10)$ gauge group, may result in conserved discrete symmetry which forbid, or suppress, their decay to the lighter Standard Model states [12].

In addition to the "Wilsonian" states the free fermionic models can give rise to semi–stable hidden glueballs that arise from the unbroken subgroup of the hidden E_8 gauge group. Such states were considered as candidates for self–interacting dark matter in ref. [8]. The hidden glueballs can interact with the Standard Model states only via super-heavy hidden matter which is additionally charged under the flavour $U(1)$ symmetries. Imposing that this hidden glueballs accounts for the dark matter requires that its self–interaction strength is of the order of hadronic interactions.

3 Candidates

3.1 The fate of the cryptons

The cryptons are fractionally charged states of the form of equation (9). This type of states appears often in realistic free fermionic models. Experimental limits on fractionally charged states impose that these states either become sufficiently massive or are confined into integrally charged states by some hidden sector gauge group. In the flipped $SU(5)$ revamped model the hidden gauge group is $SO(10) \times SU(4)$. All the fractionally charged states in this model transform under representations of the hidden $SO(10)$ or $SU(4)$ group factors. In table 1 we enumerate the spectrum of hidden matter in the revamped flipped $SU(5)$ model.

The states in the first table arise from the sectors $b_j + 2\gamma$ and $b_j + 2\gamma + \zeta$. As noted above these states transform under vectorial representations of the unbroken hidden E_8 subgroup, and are $SO(10)$ singlets. Consequently, they typically will obtain a mass term at a relatively low order in the superpotential, and can have lower level interaction terms in the superpotential. They affect the hidden sector dynamics but do not give rise to fractionally charged exotics. The states in the second table on the other hand, all transform as 4 or $\bar{4}$ of the hidden $SU(4)$ group factor, and carry fractional electric charge $\pm 1/2$. These states arise from the "Wilsonian" sectors, which are obtained from combinations of the $SO(10)$ breaking basis vector with the other basis vectors that define the revamped flipped $SU(5)$ model. Analysis of the renormalization-group β functions of $SO(10)$ and $SO(6)$ suggests that their confinement scales is of the order $\Lambda_{10} \sim 10^{14-15}$GeV for $SO(10)$ and $\Lambda_4 \sim 10^{11-12}$GeV for $SU(4)$. This indicates that the $SU(4)$ states form the lightest bound states.

The bound $SU(4)$ states are then composed of mesons, $T_i T_j$, $\Delta_i \Delta_j$ and $\tilde{F}_i \tilde{\bar{F}}_j$, baryons, $\tilde{F}_i \tilde{F}_j \Delta_k$ and $\tilde{\bar{F}}_i \tilde{\bar{F}}_j \Delta_k$, and quadrilinear *tetrons*, which are composed of four $\tilde{F}_i$s. The lightest of those have the forms $\tilde{F}_i \tilde{F}_j \tilde{F}_k \tilde{F}_l$ and $\tilde{\bar{F}}_i \tilde{\bar{F}}_j \tilde{\bar{F}}_k \tilde{\bar{F}}_l$, where $i, j, k, l = 3, 5$. As in the case of QCD pions, one may expect the charged states to be slightly

$\Delta^0_1(0,1,6,0,-\frac{1}{2},\frac{1}{2},0)$ $\Delta^0_3(0,1,6,-\frac{1}{2},-\frac{1}{2},0,\frac{1}{2})$ $\Delta^0_5(0,1,6,\frac{1}{2},0,-\frac{1}{2},0)$	$\Delta^0_2(0,1,6,-\frac{1}{2},0,\frac{1}{2},0)$ $\Delta^0_4(0,1,6,0,-\frac{1}{2},\frac{1}{2},0)$
$T^0_2(10,1,-\frac{1}{2},0,\frac{1}{2},0)$ $T^0_4(10,1,0,\frac{1}{2},-\frac{1}{2},0)$	$T^0_1(10,1,0,-\frac{1}{2},\frac{1}{2},0)$ $T^0_3(10,1,-\frac{1}{2},-\frac{1}{2},0,\frac{1}{2})$ $T^0_5(10,1,-\frac{1}{2},0,\frac{1}{2},0)$

$\tilde{F}_1^{+\frac{1}{2}}(1,4,-\frac{1}{4},\frac{1}{4},-\frac{1}{4},\frac{1}{2})$	$\tilde{F}_2^{+\frac{1}{2}}(1,4,-\frac{1}{4},\frac{1}{4},-\frac{1}{4},-\frac{1}{2})$
$\tilde{F}_3^{-\frac{1}{2}}(1,4,\frac{1}{4},\frac{1}{4},-\frac{1}{4},\frac{1}{2})$	$\tilde{F}_4^{+\frac{1}{2}}(1,4,\frac{1}{4},-\frac{1}{4},-\frac{1}{4},6-\frac{1}{2})$
$\tilde{F}_5^{+\frac{1}{2}}(1,4,-\frac{1}{4},\frac{3}{4},-\frac{1}{4},0)$	$\tilde{F}_6^{+\frac{1}{2}}(1,4,-\frac{1}{4},\frac{1}{4},-\frac{1}{4},-\frac{1}{2})$
$\tilde{\bar{F}}_1^{-\frac{1}{2}}(1,4,-\frac{1}{4},\frac{1}{4},\frac{1}{4},\frac{1}{2})$	$\tilde{\bar{F}}_2^{-\frac{1}{2}}(1,4,-\frac{1}{4},\frac{1}{4},\frac{1}{4},-\frac{1}{2})$
$\tilde{\bar{F}}_3^{+\frac{1}{2}}(1,4,-\frac{1}{4},-\frac{1}{4},\frac{1}{4},-\frac{1}{2})$	$\tilde{\bar{F}}_4^{-\frac{1}{2}}(1,4,-\frac{1}{4},\frac{1}{4},\frac{1}{4},-\frac{1}{2})$
$\tilde{\bar{F}}_5^{-\frac{1}{2}}(1,4,-\frac{3}{4},\frac{1}{4},-\frac{1}{4},0)$	$\tilde{\bar{F}}_6^{-\frac{1}{2}}(1,4,\frac{1}{4},-\frac{1}{4},\frac{1}{4},-\frac{1}{2})$

Table 1: The spectrum of hidden matter fields that are massless at the string scale in the revamped flipped $SU(5)$ model. We display the quantum numbers under the hidden gauge group $SO(10) \times SO(6) \times U(1)^4$, and subscripts indicate the electric charges.

heavier than the neutral ones, due to electromagnetic energy mass splitting. No non-renormalizable interaction capable of enabling this lightest bound state to decay has been found in a search up to eighth order. Generically, the crypton lifetime is expected to be given by

$$\tau_x \approx \frac{1}{m_x}\left(\frac{M_S}{m_x}\right)^{2(N-3)}, \tag{12}$$

where $m_x \sim \Lambda$ is the hidden sector confinement scale, M_S is the string scale and N is order of the nonrenormalizable terms that induce the tetron decay. Taking $N = 8$, $M_S \sim 10^{17-18}$ GeV and a tetron mass $m_x \sim 10^{12}$ GeV, one finds that $\tau_x > 10^{7-17}$ years. It has therefore been suggested that the lightest neutral tetron is a perfect candidate for a superheavy dark matter particle.

However, as discussed above, in addition to the lightest neutral tetron charged tetrons are formed as well. These can be argued to be slightly heavier than the neutral tetron due to the electromagnetic splitting. In order not to be over–abundant today the charged tetrons should decay into the neutral tetron. However, the charged tetron

is composed of the same constituents that compose the neutral tetron. Namely, it is composed of the constituent cryptons that carry fractional electric charge $\pm 1/2$. Suppose then that the charged tetron is composed of three $+1/2$ charged cryptons and one $-1/2$ charged crypton. In order to get a neutral tetron one of the $+1/2$ cryptons has to convert to a $-1/2$ crypton. However, since the cryptons are $SU(2)_L$ singlets, they cannot convert by emitting a charged W^+ gauge boson, or a charged Higgs. That is, the charged tetron cannot convert to a neutral one by emitting a light degree of freedom. Therefore, the only way for the charged tetron to convert to a neutral one is by exchanging a charge $+1$ heavy degree of freedom. The only possible such states are the heavy charged gauge bosons or massive string states. Effectively, therefore, the only way for the charged tetron to decay to the neutral one is by the same higher order nonrenormalizable operators which govern the decay of the neutral tetrons. In effect the reason is that both the neutral and charged tetron decay arise from annihilation, through the higher order nonrenormalizable terms, of the constituent cryptons inside the tetrons. The conclusion is therefore that if the neutral tetron is long lived, so will be the charged tetron. The abundance of the charged tetron is therefore comparable to the abundance of the neutral tetron. As stable charged matter is strongly constrained, this argument therefore indicates that the tetron cannot provide an appealing candidate for cold dark matter. One should of course qualify this statement by admitting that it is of course not impossible that specific models will contrive to give rise to operators which allow charged tetron decay while they forbid the neutral tetron decay. However, we note that in the revamped $SU(5)$ this is not the case as such operators do not arise up to order $N = 8$ nonrenormalizable terms.

To conclude this discussion, we note that the crypton passes two of the criteria that are needed to produce an appealing superheavy dark matter candidate, while it fails on the third. Namely, its mass scale is generated by the hidden sector strong dynamics and its stability arises from the "Wilsonian" symmetry breaking mechanism and the resulting fractional electric charge. However, the fact that it carries fractional electric charge implies that also the charged tetrons are long lived and give rise to the stable charged matter which is severely constrained. It is still possible, however, that the tetron exists in sufficient abundance to account for the UHECR events, while it is sufficiently diluted to evade the charged dark matter constraints. Next we examine this possibility.

3.2 Abundance of charged and neutral tetrons

Tetrons cannot have been in thermal equilibrium in the early universe, since then their freeze-out energy density ρ_T would be larger than the critical density ρ_c of the universe [25]. Hence, they must have been produced non-thermally. It has been shown that particles with masses $\gtrsim 10^{12}$ GeV can efficiently be produced gravitationally [26], even if the reheating temperature is much lower than their mass. Since

charged and neutral tetrons have almost identical masses, they should be produced in equal abundances in gravitational production mechanisms. On the other hand, heavy particles can also be produced non-thermally during preheating [27] or reheating [28]. Then the electromagnetic interactions of charged tetrons might even lead to a charged tetron abundance larger than the neutral tetron abundance, if the freeze-out temperature of charged tetrons is smaller than the maximum temperature reached during reheating. Hence, the charged tetron abundance is going to be at least of the same order as the neutral tetron abundance.

However, there are strong bounds on the abundance of long-lived charged massive particles (CHAMPs) [29]. Indeed, if neutral and charged tetrons are to constitute cold dark matter (CDM), one would expect a certain flux of CHAMPs which should be measurable, e.g. in GUT monopole detectors. Results from MACRO [30] and a surface scintillator array [31] place bounds on such a flux which are well below the expected dark matter flux [29]. Further, charged tetrons are captured in disk stars and can destroy neutron stars on a time scales $\lesssim 10$ years, i.e. for a mass $\sim 10^{12}$ GeV the tetron energy density can be at most [32]

$$\Omega_T \lesssim 10^{-6}\Omega_{CDM} \ . \tag{13}$$

Hence, tetrons cannot be cold dark matter. However, they still can be responsible for the ultrahigh energy cosmic rays if [3]:

$$\frac{\Omega_T}{\Omega_{CDM}}\frac{t_U}{\tau_T} \sim 5 \times 10^{-11} \ , \tag{14}$$

where t_U is the age of the universe and τ_T the lifetime of these tetrons. Hence, the upper bound (13) yields an upper bound on the tetron lifetime $\tau_T \lesssim 2 \times 10^4 t_U$.

3.3 Renormalization group analysis

In the previous subsections we showed that the crypton's abundance is severely limited by constraints on stable heavy charged matter, but it may still be sufficiently longed–lived and abundant to account for the UHECR events. Additionally, the cryptons affect the renormalization group equations of the Standard Model parameters. In this section we study the effect on gauge coupling unification.

In the revamped flipped $SU(5)$ model the cryptons carry fractional charge $\pm 1/2$ and transform as 4 and $\bar{4}$ under the hidden $SU(4)$ gauge group. As discussed in refs. [4, 5], from analysis of the superpotential up to order $N = 6$ it is noted that the states $\tilde{F}_{3,5}^{+\frac{1}{2}}$ and $\tilde{\bar{F}}_{3,5}^{+\frac{1}{2}}$ remain massless and confine to form the bound tetron states. These states therefore contribute to the evolution of the Renormalization Group Equations (RGEs) of the Standard Model gauge parameters. In the perturbative heterotic–string the Standard Model gauge couplings are unified at the string scale, which is of the order [36]

$$M_S \ \equiv \ M_{\rm string} \ \approx \ g_{\rm string} \times 5 \times 10^{17} \ \ {\rm GeV} \ , \tag{15}$$

where $g_{\rm string}$ is the unified string coupling. The one–loop RGEs for the Standard Model gauge couplings are given by,

$$\frac{16\pi^2}{g_i^2(\mu)} \;=\; k_i \frac{16\pi^2}{g_{\rm string}^2} \;+\; b_i \,\ln \frac{M_{\rm string}^2}{\mu^2} \;+\; \Delta_i^{\rm (total)} \tag{16}$$

where b_i are the one-loop beta-function coefficients, and the $\Delta_i^{\rm (total)}$ represent possible corrections from the additional gauge or matter states. By solving (16) for $i = 1,2,3$ simultaneously, we obtain expressions for $\sin^2\theta_W(M_Z)$ and $\alpha_3(M_Z)$, which are confronted with the experimentally measured values for these observables at the Z–boson mass scale. The expression for $\alpha_3(M_Z)$ then takes the general form

$$\begin{aligned}\alpha_3^{-1}(M_Z)|_{\overline{MS}} \;&=\; \Delta^{(\alpha)}_{\rm MSSM} + \Delta^{(\alpha)}_{\rm h.s.} + \Delta^{(\alpha)}_{\rm l.s.} + \Delta^{(\alpha)}_{\rm i.g.} + \\ &\qquad \Delta^{(\alpha)}_{\rm i.m.} + \Delta^{(\alpha)}_{\rm 2-loop} + \Delta^{(\alpha)}_{\rm Yuk.} + \Delta^{(\alpha)}_{\rm conv.},\end{aligned} \tag{17}$$

and likewise for $\sin^2\theta_W(M_Z)|_{\overline{MS}}$ with corresponding corrections $\Delta^{(\sin)}$. Here $\Delta_{\rm MSSM}$ represents the one-loop contributions from the spectrum of the Minimal Supersymmetric Standard Model (MSSM) between the unification scale and the Z scale, and the remaining Δ terms respectively correspond to the second-loop corrections, the Yukawa-coupling corrections, the corrections from scheme conversion, the heavy string thresholds, possible light SUSY threshold corrections, corrections from possible additional intermediate-scale gauge structure between the unification and Z scales, and corrections from possible extra intermediate-scale matter. Each of these Δ terms has an algebraic expression in terms of $\alpha_{\rm e.m.}$ as well as model-specific parameters such as k_1, the beta-function coefficients, and the appropriate intermediate mass scales. Here we take $k_1 = 5/3$, which is the canonical $SO(10)$ value. We also neglect here the small effect [24] of $\Delta^{(\alpha)}_{\rm h.s.}$, $\Delta^{(\alpha)}_{\rm l.s.}$, $\Delta^{(\alpha)}_{\rm 2-loop}$, $\Delta^{(\alpha)}_{\rm Yuk.}$ $\Delta^{(\alpha)}_{\rm conv.}$, as we are only interested in the qualitative effect of the intermediate crypton matter.

For $\sin^2\theta_W(M_Z)$ we have:

$$\begin{aligned}\Delta^{(\sin)}_{\rm MSSM} \;&=\; \frac{1}{k_1+1}\left[1 - \frac{a}{2\pi}(11-k_1)\ln\frac{M_S}{M_Z}\right] \\ \Delta^{(\sin)}_{\rm i.m.} \;&=\; \frac{1}{2\pi}\sum_i \frac{k_1 a}{k_1+1}\,(b_{2_i} - b_{1_i})\ln\frac{M_S}{M_i}\end{aligned} \tag{18}$$

where $M_S \equiv M_{\rm string}$ is the string unification scale, $a \equiv \alpha_{\rm e.m.}(M_Z)$, M_i are the intermediate matter mass scales. Likewise, for $\alpha_3^{-1}(M_Z)$, we have:

$$\begin{aligned}\Delta^{(\alpha)}_{\rm MSSM} \;&=\; \frac{1}{1+k_1}\left[\frac{1}{a} + \frac{1}{2\pi}(-3k_1 - 15)\ln\frac{M_S}{M_Z}\right] \\ \Delta^{(\alpha)}_{\rm i.m.} \;&=\; -\frac{1}{2\pi}\sum_i\left[\left(\frac{k_1}{1+k_1}\right)b_{1_i} + \left(\frac{1}{1+k_1}\right)b_{2_i} - b_{3_i}\right]\ln\frac{M_S}{M_i}\end{aligned} \tag{19}$$

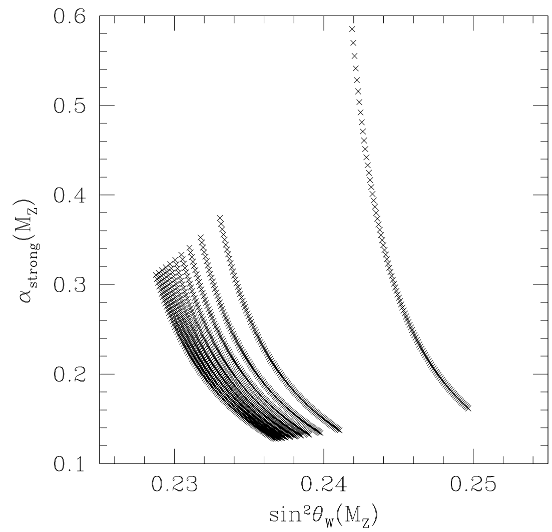

Figure 1: Scatter plot of $\sin^2\theta_W(M_Z)$ versus $\alpha_s(M_Z)$ for the entire range in eq. (21).

An important issue in string unification is that of the unification scale. The perturbative heterotic string predicts the scale of eq. (15). In this case string gauge coupling unification necessitates the existence of additional $SU(3)_{\rm colour}$ matter representations [24]. However, nonperturbative string dualities reveal that the nonperturbative string unification can be lower [37], and can be compatible with the MSSM unification scale [38] without the need for additional colour matter states. We discuss the consequences of crypton matter in both cases.

The beta–function coefficients for the cryptons $\tilde{F}$ and $\overline{\tilde{F}}$ are

$$\begin{pmatrix} b_{SU(3)} \\ b_{SU(2)} \\ b_{U(1)} \end{pmatrix}_{\tilde{F},\overline{\tilde{F}}} = \begin{pmatrix} 0 \\ 0 \\ 3/5 \end{pmatrix} . \tag{20}$$

Since we have a total of four massless representations at the scale Λ_4, the total contribution of the crypton matter states above that scale to the $U(1)$ beta–function is 12/5. We insert these numbers into the equations for the predicted values of $\alpha_3^{-1}(M_Z)$ and $\sin^2\theta_W(M_Z)$, and vary

$$\begin{aligned} 10^{12}~{\rm GeV} \le & \quad \Lambda_4 \quad \le 10^{15}~{\rm GeV} \\ 2\cdot 10^{16}~{\rm GeV} \le & \quad M_S \quad \le 5\cdot 10^{17}~{\rm GeV}. \end{aligned} \tag{21}$$

In this we in effect assume the spectrum of the MSSM plus the additional crypton states, and ignore the possible effect of the intermediate GUT threshold. The variation of the string unification scale incorporates the possible non–perturbative

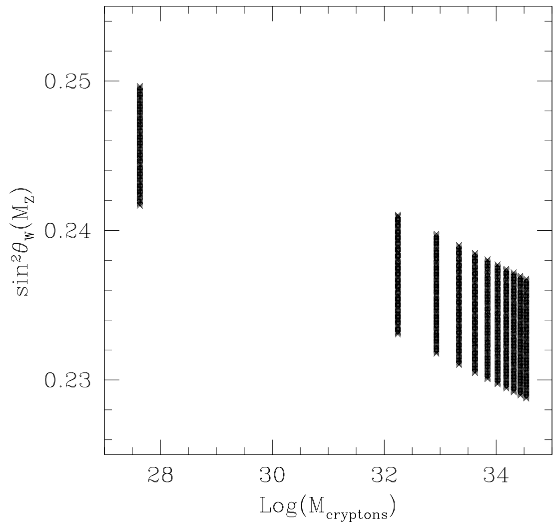

Figure 2: Scatter plot of $\sin^2\theta_W(M_Z)$ versus $\text{Log}(M_{\text{crypton}}) \equiv \Lambda_4$ for the range in eq. (21).

string effects, while that of the crypton mass scale incorporates the possible variation of the meta–stable crypton mass scale. In fig (1) we present a scatter plot of $\sin^2\theta_W(M_Z)$ versus $\alpha_s(M_Z)$ for the entire range in eq. (21). In figs. (2) and (3) we plot $\sin^2\theta_W(M_Z)$ and $\alpha_s(M_Z)$, respectively, versus $\text{Log}(M_{\text{crypton}}) \equiv \Lambda_4$. From fig. (1) we note that for the entire range in eq. (21) viable values for $\sin^2\theta_W(M_Z)$ and $\alpha_s(M_Z)$ are not obtained. This is of course somewhat expected as the additional matter affects these observables, which are in agreement with experiments if one assumes solely the MSSM spectrum in the desert. There may exist, however, a priori an interplay between the effect of additional matter and scale variation, which can result in an agreement with the experimental data. In fig. (2) we see that lower Λ_4 results in stronger disagreement with the data, whereas for $\Lambda_4 \sim 10^{15}$ GeV viable values for $\sin^2\theta_W(M_Z)$ can be obtained. Similar dependence is noted in fig. (3) for $\alpha_s(M_Z)$, with $\alpha_s(M_Z) \geq 0.127$.

In the case of the flipped $SU(5)$ model, the effect of the intermediate scale can be incorporated via the following additional term in (18):

$$-\frac{1}{2\pi}\,\frac{32}{5}\,\frac{k_1 a}{k_1+1}\,\ln\frac{M_S}{M_I}\ . \tag{22}$$

Here we have assumed the spectrum below M_I to be that of the MSSM, and above M_I to consist of three **16** representations of $SO(10)$, one **5** and $\overline{\mathbf{5}}$ of $SU(5)$ that produces the light Higgs doublets, and one **10** and $\overline{\mathbf{10}}$ of $SU(5)$ that is used to break the $SU(5)\times U(1)$ symmetry to $SU(3)\times SU(2)\times U(1)$. In addition to the variation in eq. (21), we vary the $SU(5)\times U(1)$ breaking scale between $M_x \equiv 2\cdot 10^{16}$ GeV and

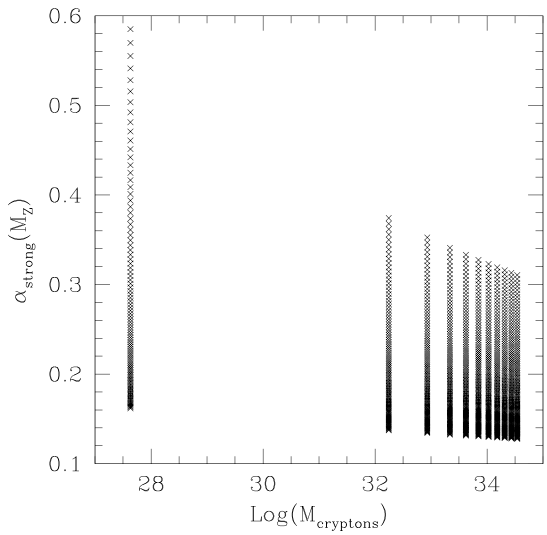

Figure 3: Scatter plot of $\alpha_s(M_Z)$ versus Log$(M_{\rm crypton}) \equiv \Lambda_4$ for the range in eq. (21).

M_s as given in eq. (21).

The existence of the intermediate $SU(5) \times U(1)$ GUT threshold does not affect the prediction for $\alpha_3^{-1}(M_Z)$, as compared to the extrapolation in the absence of this threshold. The reason being that the evolution of α_3 and α_2 is identical above this threshold. Hence, the dependence on their running above this threshold drops out when their evolution equations are equated at the unification scale. The effect of the extended gauge structure in this model is to reduce $\sin^2\theta_W(M_Z)$. In fig. (4) we present a scatter plot of $\sin^2\theta_W(M_Z)$ versus $\alpha_s(M_Z)$ for the entire range in eq. (21), including the variation of the intermediate $SU(5) \times U(1)$ GUT breaking threshold. Again we find that experimentally viable values are not obtained. In fig. (5) we show a plot of $\sin^2\theta_W(M_Z)$ versus $\alpha_s(M_Z)$ for a fixed value of $M_5 = 2 \cdot 10^{16}$ GeV and two fixed values of ($M_S = M_5; M_S = 5 \cdot 10^{17}$ GeV, and the crypton mass scale is varied as in eq. (21). The sparse (denser) curves correspond to the high (low) M_S scales, respectively. As expected the high M_S scale results in stronger disagreement with $\alpha_s(M_Z)$. However, also for the low M_S scale we see that $\sin^2\theta_W(M_Z) \geq 0.237$, and, as expected, approaches the experimentally allowed region for the larger values of Λ_4. The conclusion of this analysis is that the presence of the crypton matter states at the intermediate scale, and agreement with the experimental observables $\sin^2\theta_W(M_Z)$ and $\alpha_s(M_Z)$ necessitates the addition of additional matter beyond the MSSM. This constraint will of course be evaded if the exotic matter states are entirely neutral with respect to the Standard Model gauge group.

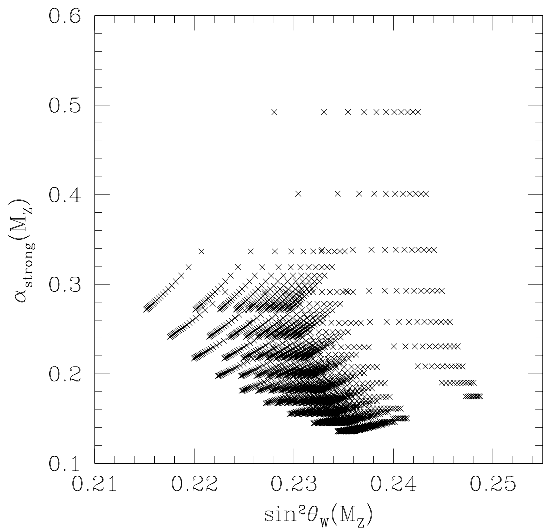

Figure 4: Scatter plot of $\sin^2\theta_W(M_Z)$ versus $\alpha_s(M_Z)$ for the entire range in eq. (21) in the presence of a $SU(5)\times U(1)$ GUT threshold.

3.4 The destiny of the uniton

The arguments above prompts us to examine whether other viable candidates for CDM and the UHECR events can arise from string models. We study the cases of the uniton, the standard model singlet and the hidden sector glueballs. In this subsection we examine whether the uniton can provide a viable candidate for the UHECR events.

The quantum numbers of the uniton are given in eq. (10). The uniton is a strongly interacting particle and carries the regular down–quark type charges under the Standard Model gauge group. It forms bound states with the light up and down quarks. The uniton carries fractional charge under the $U(1)_{Z'}$ which is embedded in $SO(10)$ and is orthogonal to the weak–hypercharge. The fractional $U(1)_{Z'}$ charge can result in local discrete symmetries that suppress its decay to the Standard Model states [12, 33]. While this is model dependent here we assume that this is indeed the case. In [12] it was argued that the uniton can give rise to a viable dark matter candidate, provided that the neutral meson is lighter than the charged one. In [34] it was shown that because of its strongly interacting light degree of freedom the uniton gets trapped in the sun and the earth in a substantial rate and subsequently annihilates into quarks and leptons. This constrains the uniton mass to be above 10^{11} GeV, in perfect agreement with the energy scale suggested by the UHECR events. The decay of the bound uniton state can then arise due to nonrenormalizable operators, yielding a lifetime similar to eq. (12).

The pertained semi–stability of the uniton arises from its fractional charge under

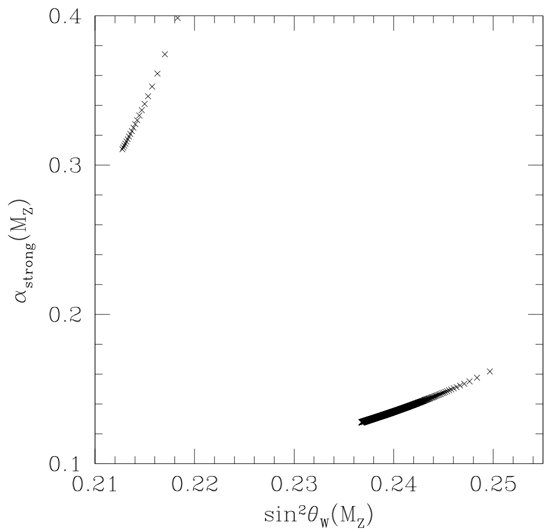

Figure 5: Scatter plot of $\sin^2\theta_W(M_Z)$ versus $\alpha_s(M_Z)$ for fixed values of M_S and M_5, and Λ_4 varied as in eq. (21).

$U(1)_{Z'}$. The neutral meson is assumed to be lighter than the charged one, in which case the uniton might provide a viable dark matter candidate. Like in the case of the crypton the question then is whether the charged meson can decay into the neutral one at a sufficient rate. The uniton itself is an $SU(2)_L$ singlet and therefore, similar to the crypton, cannot emit a light $W^\pm$ gauge boson. However, the heavy bound charged meson, contrary to the case of the tetron, is composed of the exotic uniton and a Standard Model up–type quark, which can decay through weak interactions. Therefore, the charged uniton bound state can decay into the neutral one by a Standard Model beta-decay. Hence, the uniton is a viable candidate for cold dark matter and UHECRs, as long as the neutral bound state is lighter than the charged one.

In regard to the mass scale of the uniton, unlike the crypton it does not transform under the hidden non–Abelian gauge sector. Therefore, its mass scale is not fixed by the hidden sector dynamics. The limits from trapping in the sun and earth, however, do constrain it to be in the region which is interesting from the perspective of the UHECR events. However, the mass scale is not generated dynamically and in a sense has to be put in by hand.

As has been discussed in section (3.2), meta–stable relics with masses $> 10^{11}$ GeV would overclose the universe if they had been in thermal equilibrium in the early universe. Hence, we have to consider non-thermal production mechanisms. It has been shown that super-heavy particles can efficiently be produced gravitationally [26], or during preheating [27] or reheating [28] after inflation. If we assume that meta–stable unitons form cold dark matter, their decays can explain the ultrahigh

energy cosmic rays if their lifetime is of the order (cf. eq. (14))

$$\tau_{uniton} \sim 2 \cdot 10^{10} t_U \ , \tag{23}$$

where t_U is the age of the universe.

To summarize, the proclaimed stability of the uniton arises from the "Wilsonian" gauge symmetry breaking. Unlike the case of the crypton the required mass scale does not originate from hidden sector dynamics. As the uniton appears in vector–like representations, the required mass scale can be generated by a direct superpotential mass term. Such mass terms, of the required order of magnitude can arise from nonrenormalizable terms. In this respect it is appealing that the phenomenological constraints on uniton dark matter impose $m_{\rm uniton} > 10^{11}$ GeV [34]. Lastly, the uniton is a strongly interacting particle, and one has to address experimental constraints on strongly interacting massive particles (SIMPs) as cold dark matter. Unitons will interact with ordinary matter and the cross section for their interactions with ordinary matter has been estimated to be [34]

$$\sigma \sim 10^{-26} \ {\rm cm}^2 \ . \tag{24}$$

Experimental constraints on SIMPs have recently been reanalyzed [39], and it has been shown that there are several unconstrained windows in the dark matter-proton cross section versus mass parameter space. In particular, SIMPs with masses $>$ 10^{11} GeV and cross sections with ordinary matter $< 10^{-22}$ cm^2 are viable cold dark matter candidates.

We further comment that it is of course true that the uniton also affects the evolution of the Standard Model gauge couplings. However, the original motivation to introduce the uniton at an intermediate energy scale was to enable heterotic string gauge coupling unification. Hence, the uniton can help to solve several problems at once, by enabling heterotic–string unification, by providing a substantial fraction of the dark matter, and by explaining the observed UHECR events.

3.5 The role of the singleton

The next candidate that we study is the Standard Model singlet, eq. (11), which carries fractional $U(1)_{Z'}$ charge. We refer to this state as the "singleton", and it also arises from the Wilson line breaking of the $SO(10)$ gauge symmetry. Therefore, such states can be meta–stable due to the fractional $U(1)_{Z'}$ charge, which may leave a residual local discrete symmetry after the $U(1)_{Z'}$ gauge symmetry breaking. For example, this will be the case if the states that break the $U(1)_{Z'}$ gauge symmetry carry only "integral" $U(1)_{Z'}$ charges[†]. In the model of ref. [16] some of the singleton states transform under the $SU(5)$ or $SU(3)$ hidden gauge groups, whereas others are singlets of all the non–Abelian groups in the model. One may then envision a scenario

[†]By "integral" here we mean charges that are compatible with $SO(10)$ embedding.

in which, after cancellation of the anomalous $U(1)$ D–term, the remaining light states are those that transform under a hidden non–Abelian gauge group. Then in a similar fashion to the generation of the crypton mass, the mass scale of the singleton can arise from the hidden sector dynamics. Finally the singleton is a Standard Model singlet and therefore there is no danger of producing stable charged relics. On the other hand, if we assume that the singleton transforms under a hidden gauge group, say $SU(4)$, then, similarly to the cryptons, it will form semi–stable bound states whose decay is similarly governed by the higher order nonrenormalizable operators, which do not involve $U(1)_{Z'}$ breaking VEVs, and do not violate the assumed local discrete symmetry that protect the singleton from decaying. Then, the meta–stability of the bound state arises from annihilation of the constituent singletons which is induced by the nonrenormalizable terms. We therefore conclude that the singleton states can provide appealing candidates for the UHECR events that satisfy the three desired criteria. Namely, their meta–stability originates from the Wilson line breaking effect; their mass scale may originate from hidden sector dynamics and perhaps most importantly, contrary to the crypton or the uniton, they are Standard Model singlets and therefore are electrically neutral. Finally, the non-thermal production mechanisms discussed in section (3.2) can give rise to a non-negligible singleton density, i.e., singletons are attractive cold dark matter candidates and can solve the ultrahigh energy cosmic ray problem.

3.6 Other candidates

The last candidate that we discuss are the exotic glueballs of reference [8]. These states were proposed there as candidates for strongly interacting dark matter. The basic idea is that the hidden E_8 gauge group of the heterotic string models is typically broken to a $SU(n) \times SU(m) \times U(1)^k$ subgroup, with matter states in vector–like representations. The matter states obtain intermediate scale mass terms, and the $SU(n)$ group factors condense at some scale below that, depending on the matter and gauge content. At that scale exotic glueballs form, the lightest of which is expected to be stable. The exotic glueballs can interact with the Standard Model states only through the heavy matter states, which are charged also with respect to the horizontal $U(1)$ symmetries. The flavour symmetries are broken near the string scale and therefore the exotic glueballs can interact with the Standard Model states through higher order operators, which are suppressed by the heavy mass scale. In ref. [35] it was argued that the self–interacting dark matter should have self–interactions of the order of the hadronic scale. This requirement therefore suggests that the exotic glueballs could not provide suitable candidates for the UHECR events.

4 Conclusions

The discovery of ultrahigh energy cosmic ray events with energies exceeding the GZK bound is one of the most intriguing recent experimental observations. Semi–stable superheavy matter states provide a plausible explanation for these events. In this paper we studied whether such a state could arise from string theory. Heterotic string theory gives rise to fractionally charged meta–stable Wilsonian matter states. Thus, the heterotic string, while allowing the embedding of the standard GUT structures, at the same time also produces meta–stable massive states. We discussed the various Wilsonian matter states that arise in the string models and classified them according to their charges under the unbroken $SO(10)$ subgroup at the string scale. We argued that states that carry fractional electric charge, which are confined by a hidden sector gauge group, produce semi–stable charged matter in addition to the lightest semi–stable neutral state. They are therefore severely limited by constraints on charged dark matter, and cannot give rise to a viable dark matter candidate. However, they may still be responsible for the observed UHECR events, if their lifetime is of the right order of magnitude. Further, fractionally charged matter at intermediate energy scales also affects the evolution of the Standard Model parameters and we showed that it would necessitate additional colour triplets and electroweak doublets to enable the gauge couplings to unify. While this does not exclude the cryptons as candidates for the UHECR events, it makes them, in our opinion, less attractive. We further showed that other Wilsonian matter states that carry the standard charges with respect to the Standard Model gauge group but carry fractional charges with respect to the $SO(10)$ $U(1)_{Z'}$ generator, can give rise to viable dark matter candidates that can also account for the UHECR events. The meta–stability of such states arises due to the Wilson line symmetry breaking mechanism and the fact that the Standard Model states are obtained from representations of the underlying GUT symmetry group. On the other hand the decay of such states would similarly arise from higher order nonrenormalizable operators. The other properties of these states render them more, or less, attractive. While in the case of the singleton the required mass scale could arise from hidden sector dynamics, in the case of the uniton it has to be put in by hand. However, it is very intriguing that the uniton is in fact constrained to be heavier than 10^{11} GeV [34], which is in perfect harmony with the mass scale required to explain the UHECR events. Lastly, additional coloured matter states at intermediate energy scale have been motivated from heterotic string gauge coupling unification. The singleton on the other hand provides the most likely dark matter candidate in the sense that it is a Standard Model singlet and therefore does not give rise to charged or coloured matter. Finally, forthcoming experimental data [40] and improved theoretical analysis along the lines of ref. [41] promise exciting new results with potentially ground–breaking discoveries.

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References

[1] N. Hayashida *et al*. *Phys. Rev. Lett.* **73** (1994) 3491;
D.J. Bird *et al*. *Astrophys. J.* **424** (1994) 491.

[2] K. Greisen, *Phys. Rev. Lett.* **16** (1966) 748;
G.T. Zatsepin and V.A. Kuzmin, *Pisma Zh. Eksp. Theor. Fiz.* **4** (1966) 114.

[3] V. Berezinsky, M. Kachelrieß and A. Vilenkin, Phys. Rev. Lett. **79** (1997) 4302;
V.A. Kuzmin and V.A. Rubakov, *Phys. Atom. Nucl.* **61** (1998) 1028.

[4] J. Ellis, J.L. Lopez and D.V. Nanopoulos, *Phys. Lett.* **B247** (1990) 257;
J. Ellis, G. Gelmini, J.L. Lopez, D.V. Nanopoulos and S. Sarkar, *Nucl. Phys.* **B373** (1992) 399.

[5] K. Benakli, J. Ellis, and D.V. Nanopoulos, *Phys. Rev.* **D59** (1999) 047301.

[6] M. Birkel and S. Sarkar, *Astropart. Phys.* **9** (1998) 297.

[7] I. Antoniadis, J. Ellis, J. Hagelin, and D.V. Nanopoulos, *Phys. Lett.* **B231** (1989) 65;
J. Lopez, D.V. Nanopoulos, and K. Yuan, *Nucl. Phys.* **B399** (1993) 654.

[8] A.E. Faraggi and M. Pospelov, hep-ph/0008223, to appear in Astropart. Phys.

[9] For a recent analysis see *e.g.*: D.M. Ghilencea and G.G. Ross, hep-ph/0102306.

[10] Y. Hosotani, *Phys. Lett.* **B126** (1983) 309; *Phys. Lett.* **B129** (1983) 193.

[11] X.G. Wen and E. Witten, *Nucl. Phys.* **B261** (1985) 651;
G.G. Athanasiu, J.J. Atick, Michael Dine and Willy Fischler, *Phys. Lett.* **B214** (1988) 55;
A. Schellekens, *Phys. Lett.* **B237** (1990) 363.

[12] S. Chang, C. Corianò and A.E. Faraggi, *Phys. Lett.* **B397** (1997) 76; *Nucl. Phys.* **B477** (1996) 65.

[13] A.E. Faraggi, *Phys. Rev.* **D46** (1992) 3204.

[14] A.E. Faraggi, D.V. Nanopoulos and K. Yuan, *Nucl. Phys.* **B335** (1990) 347.

[15] I. Antoniadis, G.K. Leontaris and J. Rizos, *Phys. Lett.* **B245** (1990) 161; G.K. Leontaris and J. Rizos, *Nucl. Phys.* **B554** (1999) 3.

[16] A.E. Faraggi, *Phys. Lett.* **B278** (1992) 131; *Phys. Lett.* **B274** (1992) 47.

[17] A.E. Faraggi, *Nucl. Phys.* **B387** (1992) 239; *Nucl. Phys.* **B403** (1993) 101.

[18] G.B. Cleaver, *et al*. *Phys. Lett.* **B455** (1999) 135; *Int. J. Mod. Phys.* **A16** (2001) 425; *Nucl. Phys.* **B593** (2001) 471; *Mod. Phys. Lett.* **A15** (2000) 1191; hep-ph/0002292; hep-ph/0104091.

[19] G.B Cleaver, A.E. Faraggi and C. savage, *Phys. Rev.* **D63** (2001) 066001; G.B. Cleaver, D.J. Clements, A.E. Faraggi, hep-ph/0106060.

[20] A.E. Faraggi and D.V. Nanopoulos, *Phys. Rev.* **D48** (1993) 3288; A.E. Faraggi, *Int. J. Mod. Phys.* **A14** (1999) 1663.

[21] A.E. Faraggi, *Nucl. Phys.* **B407** (1993) 57; *Phys. Lett.* **B326** (94) 62.

[22] S. Chaudhuri, G. Hockney and J. Lykken, *Nucl. Phys.* **B469** (1996) 357; G.B. Cleaver *et. al.* , *Nucl. Phys.* **B525** (1998) 3; *Nucl. Phys.* **B545** (1998) 47; *Phys. Rev.* **D59** (1999) 055005; *Phys. Rev.* **D59** (1999) 115003.

[23] S. Kalara, J. Lopez, and D.V. Nanopoulos, *Nucl. Phys.* **B353** (1991) 650.

[24] A.E. Faraggi, *Phys. Lett.* **B302** (1993) 202; K.R. Dienes and A.E. Faraggi, *Phys. Rev. Lett.* **75** (1995) 2646; *Nucl. Phys.* **B457** (1995) 409.

[25] K. Griest and M. Kamionkowski, Phys. Rev. Lett. **64** (1990) 615

[26] D. J. Chung, E. W. Kolb and A. Riotto, Phys. Rev. D **59** (1999) 023501

[27] E. W. Kolb, A. Linde and A. Riotto, Phys. Rev. Lett. **77** (1996) 4290

[28] D. J. Chung, E. W. Kolb and A. Riotto, Phys. Rev. D **60** (1999) 063504

[29] for a review and references, see M. L. Perl, P. C. Kim, V. Halyo, E. R. Lee, I. T. Lee, D. Loomba and K. S. Lackner, `hep-ex/0102033`

[30] M. Ambrosio *et al.* [MACRO Collaboration], `hep-ex/0009002`

[31] B. Barish, G. Liu and C. Lane, Phys. Rev. D **36** (1987) 2641.

[32] A. Gould, B. T. Draine, R. W. Romani and S. Nussinov, Phys. Lett. B **238** (1990) 337.

[33] A.E. Faraggi, *Phys. Lett.* **B398** (1997) 88.

[34] A.E. Faraggi, K.A. Olive and M. Pospelov, *Astropart. Phys.* **13** (2000) 31.

[35] D.N. Spergel and P.J. Steinhardt, *Phys. Rev. Lett.* **84** (2000) 3760.

[36] V.S. Kaplunovsky, *Nucl. Phys.* **B307** (88) 145; Erratum: *ibid.* **B382** (1992) 436.

[37] E. Witten, *Nucl. Phys.* **B471** (1996) 135.

[38] J. Ellis, S. Kelley, and D. V. Nanopoulos, *Phys. Lett.* **B260** (91) 131;
U. Amaldi, W. de Boer, and H. Füstenau, *Phys. Lett.* **B260** (91) 447;
P. Langacker and M. Luo, *Phys. Rev.* **D44** (91) 817;
A.E. Faraggi and B. Grinstein, *Nucl. Phys.* **B422** (1994) 3.

[39] P. C. McGuire and P. J. Steinhardt, `astro-ph/0105567`

[40] *The Pierre Auger observatory, The Auger Collaboration*, Nucl.Phys.Proc.Suppl.85 (2000) 324.

[41] C. Corianò and A.E. Faraggi, hep-ph/0106326.